Excitation of Waveguides for Integrated Optics with Laser Beams

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We discuss in this paper the excitation of dielectric waveguides used for integrated optical circuitry. Thin optical films are usually excited by means of coupling prisms using frustrated total internal reflection. We propose instead excitation of thin-film waveguides directly by gaussian laser beams. Harmful effects of imperfect substrate edges can be avoided by letting the dielectric waveguide end inside of the substrate. The exciting laser beam is directed towards the end of the guide from the inside of the substrate material.

The analysis is based on neglecting reflection at the end of the thin-film waveguide and on assuming that the thin film is infinitely extended in one dimension. The maximum excitation efficiency predicted with this model is 97 percent. It is expected that the excitation efficiency of more realistic guides for integrated optics can be as high as 90 percent.

I. INTRODUCTION

Communications systems using light waves as the carrier of information need some means of signal processing at the end terminals and perhaps also at intermediate repeater points along the transmission lines. It has been suggested[1,2] to employ integrated optical circuits for the purpose of filtering, amplifying, pulse regeneration, etc., of the optical signal beam.

These introductory remarks make it plausible that the need exists for exciting guided modes in the waveguides used for integrated optical circuits.[3] An efficient method of mode excitation utilizes the evanescent field outside of a high-index prism to couple light energy from a laser beam to one of the guided modes of a thin-film waveguide.[4] This method is particularly suitable for wide thin-film guides. However, for the very narrow light guides that are likely to be used for integrated optical circuits the prism coupler method may become less

efficient since the dielectric optical waveguide may easily be narrower than tbe laser beam. A very narrow beam must be highly focused and thus bas the disadvantage of a high-beam divergence contrary to tbe requirements of the prism coupler.

An alternate method of exciting the guided modes of dielectric waveguides is by shining tbe laser beam directly at the end of tbe guide. It has been shown theoretically that the conversion efficiency obtainable by this method is as high or higher than that of the prism coupler.[5,6] Ligbt injection by direct excitation witb a laser beam is usually employed with cladded round optical fibers. The application of the same method to the dielectric waveguides of integrated optical circuits suffers from the disadvantage that the ends of these guides are likely to be of poor optical quality. Figure 1 shows the geometry of one possible form of an integrated optical waveguide. This guide can be produced by diffusion, by sputtering or evaporation techniques.

Unless the end of the substrate of the guide is optically polished after tbe deposition of the bigher refractive index region, tbe end of tbe resulting guide can be expected to appear as shown in the figure. The optical quality of the end of the guide can be improved if tbe guide is not allowed to extend all tbe way to tbe end of tbe substrate but instead terminates inside of it as sbown in Fig. 2. Even though this method may avoid tbe problem of poor optical quality of the edge of the substrate, it introduces the new problem of how to excite the dielectric waveguide whose end is not easily accessible from the outside.

A solution to this problem is shown in Fig. 3. The laser beam is incident not from the outside of tbe dielectric substrate but from inside of it. The problem of injecting the beam into the substrate is not as severe as the problem of shining a laser beam directly on the end of the waveguide shown in Fig. 1. The beam enters the substrate

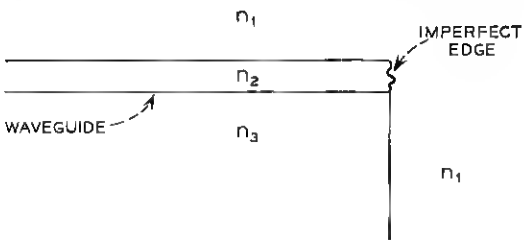

Fig. 1—Thin-film waveguide located directly under the surface of the substrate extending all the way to its edge.

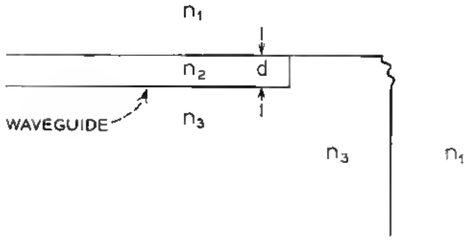

Fig. 2—Thin-film waveguide not reaching to the edge of the substrate.

not at its edge but in a region that is more easily kept free of damage. An external focusing system is, of course, required to achieve a beam of the desired convergence angle.

To study the required beam pattern for the efficient excitation of the optical waveguide, we reverse the problem and study instead the radiation pattern produced by a guided mode reaching the end of the waveguide of Fig. 3. The shape of the laser beam used for mode excitation must be made as similar as possible to the radiation pattern of the waveguide with the only exception that the direction of travel of the light field be reversed.

In the bulk of the paper we discuss the far field radiation pattern of a guided mode leaving the end of the waveguide inside of the substrate material. For simplicity we assume that the waveguide has the form of a slab that is infinitely extended in a direction perpendicular to the plane of Fig. 3. The excitation problem of the actual waveguide is very similar to that of the slab. In particular it is well known that a mode can be excited with high efficiency by an external laser beam. The only complication introduced into the problem under discussion is the presence of the air-dielectric interface of the sub-

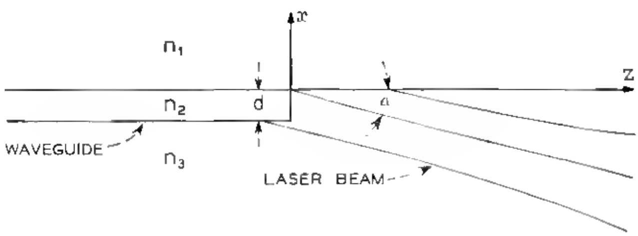

Fig. 3—Thin-film waveguide geometry used for the calculation.

strate material. This complication exists only in the direction parallel to the plane of the drawing. The mode matching problem in the direction perpendicular to the plane of the figure is no more difficult than that treated in earlier publications.[6]

A brief discussion of sligbtly different waveguide geometries can be found in the section on conclusions at the end of the paper.

II. MODES OF THE DIELECTRIC MEDIUM

The radiation pattern of the guided mode leaving the end of the dielectric waveguide could be computed with the conventional methods of the Kirchhoff-Huygens diffraction integral if it were not for the presence of the air-dielectric interface of the substrate material. A substantial part of tbe radiation field is reflected at this interface complicating the beam pattern. It appears more natural to employ the normal mode method for the solution of the radiation problem. This method utilizes the fact that any field can be expressed as a superposition of a complete set of normal modes. The air-dielectric interface is thus included automatically. We assume that the difference between tbe refractive index of the waveguide and the substrate is so slight that reflection at the end of the waveguide can be neglected. The radiation field is then obtained by the requirement that the transverse electric field component must be continuous in the plane $z = 0$ of Fig. 3. Tbe radiation field is expressed as a superposition of the normal modes of the structure to the right of the end of the waveguide. No guided modes exist in this region so that the superposition of modes must be expressed as an integral over the continuum of radiation modes.

Since we have simplified the problem to that of a slab infinitely extended in y direction, we impose the condition that there is no field variation in that direction

$$\frac{\partial}{\partial y} = 0. \tag{1}$$

The guided mode field is assumed to be of the form

$$\left.\begin{aligned} E_y &= A_e \exp(-\gamma x) && 0 \leqq x \leqq \infty \\ E_y &= A_e\left(\cos \kappa x - \frac{\gamma}{\kappa}\sin \kappa x\right) && -d \leqq x \leqq 0 \\ E_y &= A_e\left(\cos \kappa d + \frac{\gamma}{\kappa}\sin \kappa d\right)\exp\left[\theta(x + d)\right] && -\infty \leqq x \leqq -d \end{aligned}\right\} z \leqq 0 \tag{2}$$

with

$$A_e = 2\kappa\left\{\frac{\omega\mu_0 P}{\beta_0\left(d + \frac{1}{\gamma} + \frac{1}{\theta}\right)(\kappa^2 + \gamma^2)}\right\}^{\frac{1}{2}}. \tag{3}$$

A factor $\exp\,[i(\omega t - \beta_0 z)]$ has been suppressed. The constants and parameters appearing in these equations have the following meaning:

ω = $2\pi f$ radian frequency,
P = power carried by the mode,
μ_0 = magnetic susceptibility of free space,
β_0 = propogation constant,
d = slab thickness,

$$\kappa = (n_2^2 k^2 - \beta_0^2)^{\frac{1}{2}}, \tag{4}$$

$$\gamma = (\beta_0^2 - n_1^2 k^2)^{\frac{1}{2}}, \tag{5}$$

$$\theta = (\beta_0^2 - n_3^2 k^2)^{\frac{1}{2}}, \tag{6}$$

k = $\omega(\epsilon_0\mu_0)^{1/2}$ free space propogation constant,
n_1 = refrective index in the region $x > 0$,
n_2 = refractive index in the region $-d < x < 0$,
n_3 = refractive index in the region $-\infty < x < -d$.

The magnetic field components follow from the relations

$$H_x = \frac{-i}{\omega\mu_0}\frac{\partial E_y}{\partial z}, \tag{7}$$

$$H_z = \frac{i}{\omega\mu_0}\frac{\partial E_y}{\partial x}. \tag{8}$$

The values of the propagation constant β_0 are obtained from the eigenvalue equation

$$\tan \kappa d = \frac{\gamma + \theta}{\kappa\left(1 - \frac{\gamma\theta}{\kappa^2}\right)}. \tag{9}$$

In addition to the guided modes there is also a continuum of radiation modes in the region $z < 0$. However, these modes are of no interest to us.

To tbe rigbt of the region where the dielectric waveguide has ended ($z > 0$) no guided modes are possible. The continuum of radiation modes can be subdivided into three distinct groups. We use the propa-

gation constant β of the continuum modes for their classification. In the region

$$0 \leqq \beta < n_1 k \tag{10}$$

the following radiation modes exist:

$$\mathcal{E}_y^{(1)} = \begin{cases} A_r \cos \rho x & 0 \leqq x < \infty, \\ A_r \cos \sigma x & -\infty < x \leqq 0, \end{cases} \tag{11}$$

and

$$\mathcal{E}_y^{(2)} = \begin{cases} iA_r \sqrt{\dfrac{\sigma}{\rho}} \sin \rho x & 0 \leqq x < \infty, \\ iA_r \sqrt{\dfrac{\rho}{\sigma}} \sin \sigma x & -\infty < x \leqq 0. \end{cases} \tag{12}$$

A factor exp $[i(\omega t - \beta z)]$ has again heen suppressed. The amplitude A_r is related to power hy the relation

$$A_r = 2\left(\frac{\omega\mu\rho P}{\pi\beta(\sigma + \rho)}\right)^{\frac{1}{2}}. \tag{13}$$

However, it must he pointed out that the power carried hy one individual radiation mode is infinite. The power P appearing in equation (13) is defined hy the expression

$$P\delta(\rho - \rho')\delta_{\nu\nu'} = \frac{\beta}{2\omega\mu} \int_{-\infty}^{\infty} \mathcal{E}_y^{(\nu)*}(\rho)\, \mathcal{E}_y^{(\nu')}(\rho')\, dx. \tag{14}$$

[$\delta(\rho - \rho')$ is Dirac's delta function, $\delta_{\nu\nu'}$, is the Kronecker delta symhol.] Equation (14) expresses the orthogonality of the radiation modes. It implies that the modes (11) and (12) are not only orthogonal among each other hut also that the modes given by (11) are orthogonal to the modes given by equation (12). The parameters ρ and σ are related to the propagation constant β by the following expressions

$$\rho = (n_1^2 k^2 - \beta^2)^{\frac{1}{2}} \tag{15}$$

and

$$\sigma = (n_3^2 k^2 - \beta^2)^{\frac{1}{2}}. \tag{16}$$

Whereas there are two sets of radiation modes in the interval given hy equation (10) there is only one set of radiation modes in the interval

$$n_1 k \leqq \beta \leqq n_3 k. \tag{17}$$

This set is given by the equation

$$\mathcal{E}_y^{(3)} = \begin{cases} B_r \exp(-\delta x) & 0 \leqq x < \infty, \\ B_r\left(\cos \sigma x - \dfrac{\delta}{\sigma} \sin \sigma x\right) & -\infty < x \leqq 0, \end{cases} \tag{18}$$

with

$$B_r = \left(\frac{4\omega\mu\sigma\delta P}{\pi\beta(\sigma^2 + \delta^2)}\right)^{\frac{1}{2}}. \tag{19}$$

The parameter σ is given by equation (16) while δ is defined as

$$\delta = (\beta^2 - n_1^2k^2)^{\frac{1}{2}}. \tag{20}$$

The modes (18) are orthogonal among each other as well as to the modes (11) and (12).

Actually the radiation modes (11) and (12) cover also the range along the imaginary β axis from 0 to $i\,\infty$. However, the modes belonging to this imaginary branch of β values do not carry power and are thus of no interest to our investigation.

A more detailed study reveals that the radiation modes of equations (11) and (12) do not take part in forming the main lobe of the radiation field in the space $z > 0$. The main lobe is obtained by a superposition of the radiation modes (18).

III. CALCULATION OF THE RADIATION PATTERN

We simplify the calculation of the radiation pattern (resulting from the guided mode carried in the waveguide at $z < 0$ spilling over into the half space $z > 0$) by matching the transverse component E_y of the electric field at the plane $z = 0$. The continuity condition for the transverse **H** component is ignored. This approximation is acceptable if reflection at the plane $z = 0$ can safely be neglected.

The continuity condition for the transverse **E** field can be expressed by the equation

$$E_y = \int_0^{(n_2^2 - n_1^2)^{\frac{1}{2}}k} q(\delta)\, \mathcal{E}_y^{(3)}(x, \delta)\, d\delta. \tag{21}$$

The field on the left side of this equation is the guided mode at $z = 0$ while the field on the right side is the radiation field at $z = 0$ expressed as a superposition of radiation modes. Only the main radiation lobe is included by restricting the mode expansion to the modes of

equation (18). The much smaller side lobes are, however, of no interest to us.

Using the orthogonality of the radiation modes we obtain the expansion coefficient $q(\delta)$ from equation (21)

$$q(\delta) = \frac{\beta}{2\omega\mu P}\int_{-\infty}^{\infty} E_y(x)\, \mathcal{E}_y^{(3)}(x,\, \delta)\, dx = 2\kappa(\beta\sigma\delta)^{1/2}(n_2^2 - n_3^2)k^2$$
$$\cdot \frac{(\gamma - \delta)(\kappa^2 + \delta^2)^{1/2} + (\kappa^2 + \gamma^2)^{1/2}\left[(\theta + \delta)\cos\sigma d + \left(\frac{\theta\delta}{\sigma} - \sigma\right)\sin\sigma d\right]}{\left[\pi\beta_0\left(d + \frac{1}{\gamma} + \frac{1}{\theta}\right)(\sigma^2 + \delta^2)(\kappa^2 + \gamma^2)(\kappa^2 + \theta^2)\right]^{1/2}(\kappa^2 - \sigma^2)(\gamma^2 - \delta^2)}. \tag{22}$$

In principle the radiation problem is thus solved. Substitution of equations (22) and (18) into (21) (restoring the omitted exponential z dependence) allows us to calculate the main lobe of the radiation field at any point $z > 0$. An exact evaluation of the integral is not possible. However, a good far field approximation can be obtained with the method of stationary phase. The result of this far field approximation is ($x < 0, z > 0$)

$$E_y = \frac{xz}{r^{5/2}} \frac{2^{3/2}\kappa(\omega\mu P)^{1/2}(n_2^2 - n_3^2)n_3^{3/2}k^{7/2}\exp\left\{i\left[-n_3kr + \arctan\frac{\delta}{\sigma}\right]\right\}}{\left[\pi\beta_0\left(d + \frac{1}{\gamma} + \frac{1}{\theta}\right)(\kappa^2 + \gamma^2)(\kappa^2 + \theta^2)(\sigma^2 + \delta^2)\right]^{1/2}(\kappa^2 - \sigma^2)(\gamma^2 - \delta^2)}$$
$$\cdot\Big\{(\gamma - \delta)(\kappa^2 + \theta^2)^{1/2} +$$
$$(\kappa^2 + \gamma^2)^{1/2}\cdot\left[(\theta + \delta)\cos\sigma d + \left(\frac{\theta\delta}{\sigma} - \sigma\right)\sin\sigma d\right]\Big\}. \tag{23}$$

The parameters entering this equation are now related to the x, z coordinates

$$\left.\begin{aligned} \beta &= n_3k\frac{z}{r} \\ \sigma &= -n_3k\frac{x}{r} \\ \delta &= (n_3^2z^2 - n_1^2r^2)^{1/2}\frac{k}{r} \\ r &= (x^2 + z^2)^{1/2} \end{aligned}\right\}. \tag{24}$$

Equation (23) is valid only in the region where δ is real. Equation (23) does not have any poles since the numerator vanishes simultaneously with the denominator.

Figure 4 shows plots of the absolute value of E_y as a function of the angle

$$\alpha = \arctan \frac{|x|}{z} \tag{25}$$

for the case $n_1 = 1$, $n_2 = 1.51$ and $n_3 = 1.5$ and for several values of kd. The most interesting aspect of these curves is their width and the position of their maxima. It is easiest to visualize the meaning of these curves by assuming that the wavelength is fixed but that the guide width d is changing. For small values of d tbe guided mode field extends far into the substrate medium with refractive index n_3. A wide aperture field leads to a narrow radiation lobe. This explains the narrow radiation pattern for the curve with $kd = 9$. As the width d increases the guided mode contracts at first giving rise to a wider radiation pattern. At $kd = 20$ the radiation pattern assumes its greatest width. As d increases even more the guided mode field again becomes wider and thus creates a narrower radiation lobe.

The position of the maximum is shown as the solid line in Fig. 5. A qualitative explanation for the position of the radiation maximum can be given as follows. The guided mode field inside of the medium

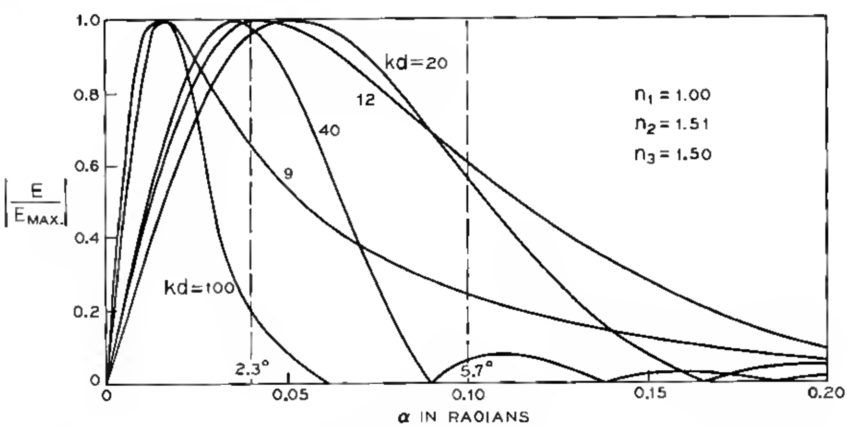

Fig. 4—Ratio of E_y over the peak value of the electric vector of the radiation field as a function of the angle in radians. $n_1 = 1.00$, $n_2 = 1.51$, $n_3 = 1.50$.

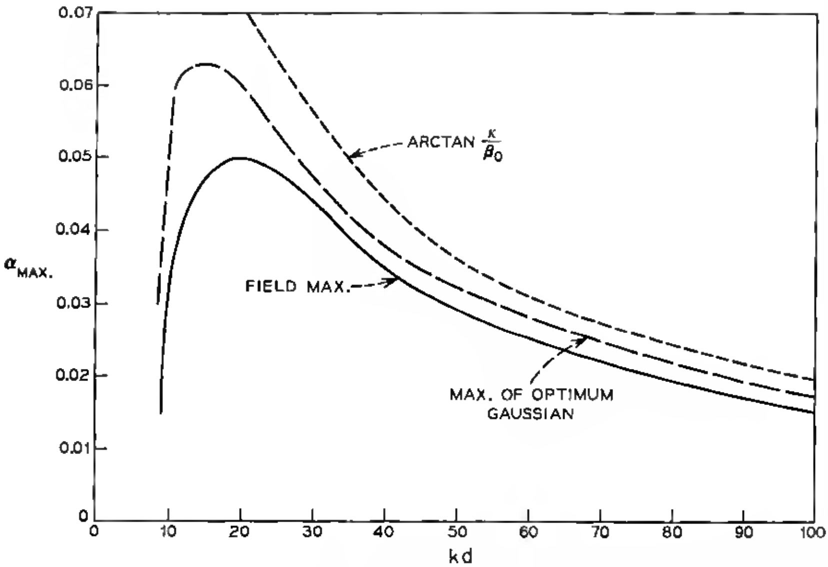

Fig. 5—Location of the peak of the radiation pattern. The solid curve indicates the maxima of the curves of Fig. 4; the dotted curve represents the direction of travel of the plane waves whose superposition comprise the guided mode; the dash-dotted curve indicates the location of the field maxima of the optimum gaussian distribution used to excite the waveguide.

with index n_2 can he decomposed into two plane waves traveling at certain slant angles given hy the equation

$$\tan \alpha' = \frac{\kappa}{\beta_0}. \tag{26}$$

It might be expected that the radiation field can he explained hy assuming that the two plane waves composing the guided mode simply detach themselves from the guidance structure and travel out into the space $z > 0$. The wave whose slope points up towards the air-dielectric interface is totally reflected and is forced to travel in the same direction as the other plane wave component. The dotted curve of Fig. 5 shows the angle α' of equation (26) as a function of kd. The dotted and solid curves do not coincide indicating that the picture of the plane waves simply detaching themselves to form the radiation pattern is not quite accurate. However, it is also apparent that this explanation provides a first rough idea of the direction into which the radiation is emitted.

Figure 6 is a plot of the phase angle

$$\phi = \arctan \frac{\delta}{\sigma} \tag{27}$$

appearing in equation (23) as a function α. This curve is independent of kd and holds for all radiation patterns shown in Fig. 4. It is apparent that the surface $r =$ const. is not exactly a surface of constant phase. Since the departure from constant phase is linear, Fig. 6 indicates that the radiation pattern does not originate at the coordinate origin but at a slightly shifted point. The shift of the far field source point is so slight however (only a few percent of the width of the guiding region) that its exact position is not important. For all practical purposes it can be assumed that the source point of the far field radiation pattern is located at the air-dielectric interface at the point where the guiding structure ends.

IV. MATCHING THE RADIATION PATTERN WITH A GAUSSIAN BEAM

We are now in a position to return to the primary purpose of our investigation. Knowing the radiation pattern produced by a guided mode we can also answer the question: what is the efficiency with which this mode can be excited? To do this we need to find what fraction of the radiation pattern appears in form of a gaussian beam. The efficiency with which the radiation field excites a given gaussian beam mode of free space is, by the reciprocity theorem, identical to the efficiency with which the radiation field and consequently the

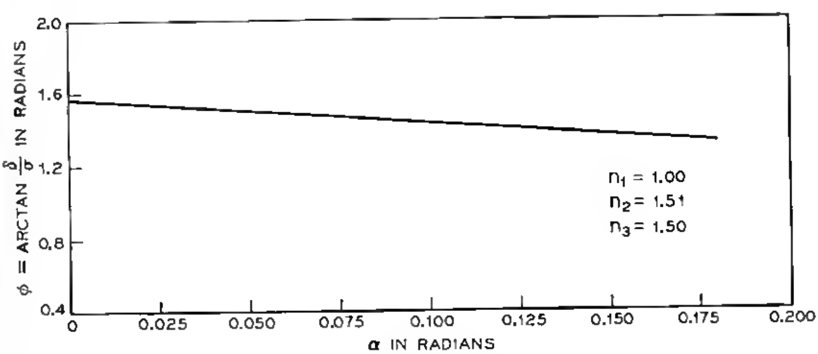

Fig. 6—The phase angle occurring in equation (23) as a function of α.

guided mode producing this radiation field can be excited by a gaussian beam mode.

The radiation field E_y can be expressed as a sum of hermite-gaussian modes $\mathcal{E}_n$

$$E_y = \sum_{n=0}^{\infty} c_n \mathcal{E}_n \, . \tag{28}$$

The only coefficient of interest to us is c_0, the excitation coefficient of the lowest order hermite-gaussian mode. Since the zero order hermite polynominal is the constant unity, this mode is purely gaussian. Using the gaussian mode

$$\mathcal{E}_0 = \left(\frac{2}{\pi}\right)^{\frac{1}{4}} \left(\frac{2\omega\mu P}{n_3 kr\theta_0}\right)^{\frac{1}{2}} \exp\left[-\left(\frac{\alpha - \alpha_0}{\theta_0}\right)\right] \tag{29}$$

with the angular beam half width θ_0 we can express the expansion coefficient c_0 as follows.

$$c_0 = \frac{n_3 k}{2\omega\mu P} \int_0^{\alpha_c} E_y \mathcal{E}_0^* r \, d\alpha$$

$$= \frac{2^{5/4} n_3^2 (n_2^2 - n_3^2) \kappa k^4}{\pi^{3/4} \left[\theta_0 \beta_0 \left(d + \frac{1}{\gamma} + \frac{1}{\theta}\right)(\kappa^2 + \gamma^2)(\kappa^2 + \theta^2)\right]^{1/2}}$$

$$\cdot \int_0^{\alpha_c} \exp\left[-\left(\frac{\alpha - \alpha_0}{\theta_0}\right)^2\right] \frac{\sin\alpha \cos\alpha}{(\sigma^2 + \delta^2)^{1/2}(\kappa^2 - \sigma^2)(\gamma^2 - \delta^2)}$$

$$\cdot \left\{ (\gamma - \delta)(\kappa^2 + \theta^2)^{1/2} + (\kappa^2 + \gamma^2)^{1/2} \right.$$

$$\left. \cdot \left[(\theta + \delta) \cos \sigma d + \left(\frac{\theta\delta}{\sigma} - \sigma\right) \sin \sigma d \right] \right\} d\alpha. \tag{30}$$

The upper limit α_c of the integral corresponds to the angle at which δ of equation (24) vanishes. The transmission coefficient from the free space gaussian beam mode to the guided mode of the waveguide is now given by

$$T = |c_0|^2. \tag{31}$$

The integral in (30) was evaluated numerically. The results are shown in Fig. 7. The optimum angles for the maximum of the gaussian beam α_0 and the optimum half width θ_0 had to be determined by trial and error. The optimum values α_0 of the peak of the gaussian beam distribution are shown as the dash-dotted curve of Fig. 5. This curve

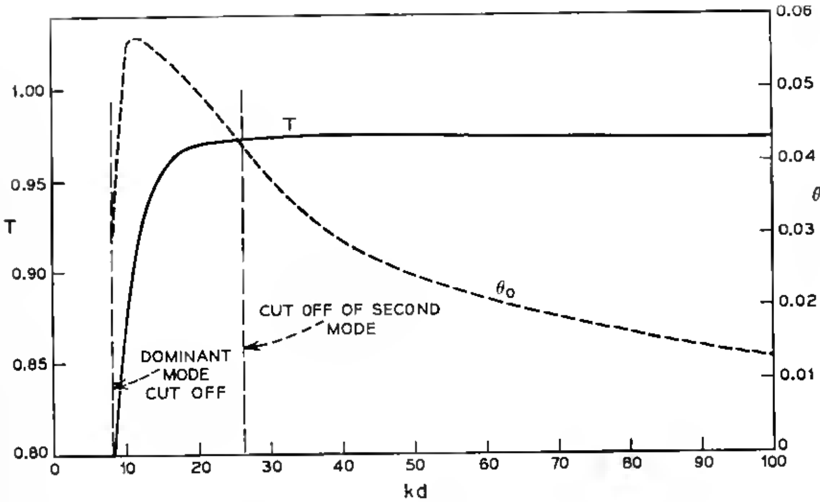

Fig. 7—Optimum transmission (excitation) coefficient T and optimum beam half width θ_0 of the exciting gaussian beam.

falls between the position of the maxima of the radiation pattern and the dotted line indicating the direction of the plane waves composing the guided mode. The optimum position of the gaussian beam mode approaches the dotted curve more closely. The idea of determining the required direction of the incident laser beam from the direction of the plane waves comprising the guided mode is thus more accurately reflected in the maximum position of tbe optimum gaussian beam.

The optimum half width θ_0 of the gaussian distribution is also shown in Fig. 7. The optimum position of tbe gaussian beam as well as its optimum width can not be obtained simply from the curves of Fig. 4 because of the asymmetry present in tbese curves.

The maximum achievable coupling efficiency is quite high. According to Fig. 7, 97 percent of the power can be converted from an optimally placed gaussian beam to the guided mode of the waveguide. We must remember, however, that our present case was oversimplified by assuming that the dielectric waveguide is infinitely extended in y direction. If the waveguide has only a finite width in y direction the gaussian beam must also be focused in that direction in order to match it to the guided mode. However, the matching problem in the y direction is not complicated by a dielectric interface. The radiation pattern that results from the finite beam width in y direction is perfectly

symmetric. Matching to symmetric radiation fields has been discussed in earlier work[5,6] and was found to be highly efficient. (99 percent of the laser power can theoretically be converted to the dominant mode of a round optical fiber.) The loss in excitation efficiency caused by a finite width of the waveguide in y direction is expected to be slight. It is safe to assume that 90 percent conversion efficiency can be obtained for a more realistic waveguide with finite widtb in x as well as y direction.

V. CONCLUSION

We have seen that a dielectric optical waveguide constructed as shown in Fig. 2 can be excited by a gaussian laser beam with high efficiency. Based on the assumption that the waveguide is infinitely extended in y direction tbe maximum obtainable excitation efficiency is 97 percent. For a more realistic guide whose y dimension is also finite tbe conversion efficiency must be somewhat less. However, based on experience with mode excitation by gaussian beams in symmetrical waveguides we expect the excitation efficiency to be better than 90 percent.

The gaussian beam must be inserted into the waveguide from the side of the substrate. This requires shining the beam into tbe substrate material at a small angle witb respect to its air-dielectric interface. A typical angle is of tbe order of 2°. It may be necessary to sbape the substrate in order to facilitate the injection of the laser beam. The main advantage of this scheme over injection of the beam into a waveguide that reaches to the edge of the substrate consists in avoiding tbe imperfections in geometry that must be expected to exist at tbe edge of the substrate material. The waveguide extends only a few microns into the substrate material so tbat the edge of the substrate would have to be absolutely perfect to within a fraction of a micron.

Other waveguide geometries can be excited by a similar tecbnique. It is possible to create a region of higher refractive index inside of the substrate material by ion implantation using high energy ion beams. A waveguide of this kind is shown schematically in Fig. 8. The waveguide may reach the outer face of tbe substrate material or it may end inside of the substrate. Excitation of this guide witb a gaussian laser beam is possible in eitber case. This geometry is particularly advantageous since the radiation pattern of this type of waveguide is

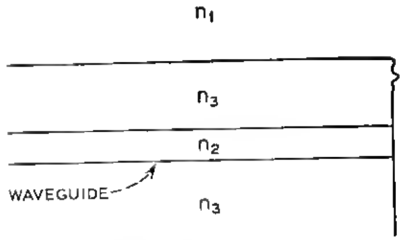

Fig. 8—Alternate dielectric waveguide configuration. The guiding region is located beneath the surface of the substrate.

not influenced by the presence of the air-dielectric interface. High excitation efficiency can be expected.

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